

# Nuclear dependence of azimuthal asymmetry in semi-inclusive deep inelastic scattering

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Within the framework of a generalized factorization, semi-inclusive deeply inelastic scattering (SIDIS) cross sections can be expressed as a series of products of collinear hard parts and transverse-momentum-dependent (TMD) parton distributions and correlations. The azimuthal asymmetry  $\langle \cos \phi \rangle$  of unpolarized SIDIS in the small transverse momentum region will depend on both twist-2 and 3 TMD quark distributions in target nucleons or nuclei. Nuclear broadening of these twist-2 and 3 quark distributions due to final-state multiple scattering in nuclei is investigated and the nuclear dependence of the azimuthal asymmetry  $\langle \cos \phi \rangle$  is studied. It is shown that the azimuthal asymmetry is suppressed by multiple parton scattering and the transverse momentum dependence of the suppression depends on the relative shape of the twist-2 and 3 quark distributions in the nucleon. A Gaussian ansatz for TMD twist-2 and 3 quark distributions in nucleon is used to demonstrate the nuclear dependence of the azimuthal asymmetry and to estimate the smearing effect due to fragmentation.

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## I. INTRODUCTION

In high-energy nuclear collisions, from deeply inelastic scattering (DIS) off a nuclear target to hadron-nucleus and nucleus-nucleus collisions, multiple parton scattering plays an important role in the reaction dynamics and the final hadron spectra. It not only causes the transverse momentum broadening of the propagating parton but also leads to parton energy loss due to induced gluon bremsstrahlung. A direct consequence of the transverse momentum broadening is the so-called Cronin effect or the broadening of the final hadron spectra in transverse momentum in both DIS off nuclear targets [1, 2] and hadron-nucleus collisions [3, 4]. Gluon bremsstrahlung induced by multiple parton scattering on the other hand leads to parton energy loss [5–8] and medium modification of the parton fragmentation functions [9, 10], a phenomenon known as jet quenching which is observed as the suppression of the leading hadron yields from parton fragmentation. The radiative parton energy loss and medium modification to parton fragmentation functions are also determined by the jet transport parameter or the average squared transverse momentum broadening per unit distance experienced by the propagating parton, which is related to the gluonic structure of the scattering centers in the medium [6, 11]. Therefore, extraction of the jet transport parameter through experimental measurements of transverse momentum broadening and jet quenching can provide important information about the properties of the medium, either cold nuclei or hot and dense QCD matter, as probed by energetic partons. The jet quenching phenomenon has been seen in high-energy heavy-ion collisions via strong suppression of not only the large transverse momentum single hadron spectra [12, 13] but also the back-to-back dihadron correlation [14]. The extracted jet transport parameter is found to be much larger than that in large cold nuclei as obtained by phenomenological studies of jet quenching in DIS [15, 16], indicating much higher gluonic density in the initial stage of high-energy heavy-ion collisions.

Much efforts have also been devoted to the study of transverse momentum broadening in DIS off nuclear targets [6, 9, 10, 17, 18, 20–24] and the Drell-Yan dilepton production in proton-nucleus collisions [25–27] within different approaches such as the color dipole model [17–19] and higher-twist expansion in the generalized collinear factorization formalism [20, 21]. It has been shown recently [24] that the gauge invariant transverse momentum dependent (TMD) quark distributions in nucleons and nuclei can be expressed as a sum of higher-twist collinear parton matrix elements,

$$f_q^A(x, k_\perp) = \int \frac{dy^-}{2\pi} e^{ixp^+ y^-} \langle A | \bar{\psi}(0) \frac{\gamma^+}{2} e^{\vec{W}_\perp(y^-) \cdot \vec{\nabla}_{k_\perp}} \psi(y^-) | A \rangle \delta^{(2)}(\vec{k}_\perp), \quad (1)$$

in terms of the parton transport operator

$$\vec{W}_\perp(y^-) \equiv i\vec{D}_\perp(y^-) + g \int_{-\infty}^{y^-} d\xi^- \vec{F}_{+\perp}(\xi^-), \quad (2)$$

where  $\vec{D}_\perp(y^-) = \vec{\partial}_\perp + ig\vec{A}_\perp(y^-)$  is the covariant derivative. For brevity of presentation, we have used the light-cone gauge and all the transverse coordinates have been set to  $\vec{y}_\perp = \vec{0}_\perp$  in the above. Since the above expression is valid

for both nucleon and nuclear targets, the nuclear dependence of TMD parton distribution functions will come from the nuclear dependence of the higher-twist parton matrix elements. Under the “maximal two gluon approximation”, the TMD quark distribution in a nucleus can be expressed as a convolution of the TMD quark distribution  $f_q^N(x, \vec{\ell}_\perp)$  in nucleon and a Gaussian broadening [24],

$$f_q^A(x, k_\perp) \approx \frac{A}{\pi \Delta_{2F}} \int d^2 \ell_\perp e^{-(\vec{k}_\perp - \vec{\ell}_\perp)^2 / \Delta_{2F}} f_q^N(x, \ell_\perp). \quad (3)$$

The broadening width  $\Delta_{2F}$  or the total average squared transverse momentum broadening,

$$\Delta_{2F} = \int d\xi_N^- \hat{q}_F(\xi_N). \quad (4)$$

is given by the quark transport parameter,

$$\hat{q}_F(\xi_N) = \frac{2\pi^2 \alpha_s}{N_c} \rho_N^A(\xi_N) [x f_g^N(x)]_{x=0}, \quad (5)$$

where  $\rho_N^A(\xi_N)$  is the spatial nucleon number density inside the nucleus and  $f_g^N(x)$  is the gluon distribution function in a nucleon,

$$x f_g^N(x) = - \int \frac{d\xi^-}{2\pi p^+} e^{ixp^+ \xi^-} \langle N | F_{+\sigma}(0) F_+{}^\sigma(\xi^-) | N \rangle. \quad (6)$$

A summation over the gluon’s color index is implied in the above definition of the gluon distribution function.

One can generalize the above approach to nuclear modification of higher-twist TMD parton distributions. In this paper, we will study the case of twist-3 TMD quark distributions which determines the azimuthal asymmetry  $\langle \cos \phi \rangle$  of the unpolarized semi-inclusive deep-inelastic scattering (SIDIS) cross section defined with respect to the leptonic plane. Such asymmetry at large transverse momentum arises predominately from hard gluon bremsstrahlung in perturbative QCD [28, 29]. Similar processes are also responsible for the azimuthal angle dependence of the Drell-Yan dilepton production cross sections [30] and multiple parton scattering in  $p + A$  collisions can lead to a twist-4 nuclear dependence [31, 32]. In the small transverse momentum  $k_\perp \leq 1$  GeV/c region, however, the asymmetry is shown to arise mainly from the intrinsic transverse momentum of the quarks in the target nucleon or nucleus [33, 34] and is related to the higher-twist TMD parton distribution functions [35–38]. We will focus on the twist-3 contribution to the azimuthal dependence, proportional to  $(k_\perp/Q) \cos \phi$ , in this paper. Here  $Q^2 = -q^2$  and  $q$  is the four momentum transfer from lepton to target. Since the coefficient of the azimuthal asymmetry is related to twist-3 TMD parton distributions, one can take the same approach as in the case of the twist-2 TMD parton distribution to study the nuclear dependence of the twist-3 contribution to the azimuthal asymmetry of SIDIS cross sections.

The remainder of this paper is organized as follows. In Sec. II, we will review the calculation of the azimuthal dependence of SIDIS up to twist-3 contributions in terms of the TMD parton distributions. In Sec. III, we extend the study of nuclear dependence of TMD parton distributions to the case of other TMD parton correlation functions, including twist-3. Using these results, we calculate the general nuclear dependence of  $\langle \cos \phi \rangle$ . We will then illustrate the nuclear dependence with an ansatz of the TMD parton distributions in a Gaussian form and discuss the effect of transverse momentum smearing in the fragmentation. A summary will be given in Sec. IV.

## II. AZIMUTHAL ASYMMETRY IN SIDIS

We consider semi-inclusive deep inelastic scattering  $e^- + A \rightarrow e^- + q + X$  with unpolarized beam and nucleus or nucleon target. The differential cross section,

$$d\sigma = \frac{\alpha_{em}^2 e_q^2}{s Q^4} L^{\mu\nu}(l, l') \frac{d^2 W_{\mu\nu}}{d^2 k'_\perp} \frac{d^3 l' d^2 k'_\perp}{(2E_{l'})}, \quad (7)$$

can be expressed as a product of the leptonic tensor,

$$L^{\mu\nu}(l, l') = 4[l^\mu l'^\nu + l^\nu l'^\mu - (l \cdot l') g^{\mu\nu}], \quad (8)$$

and the hadronic tensor,

$$\frac{d^2 W_{\mu\nu}}{d^2 k'_\perp} = \int \frac{dk'_z}{(2\pi)^3 2E_{k'}} W_{\mu\nu}^{(si)}(q, p, k'); \quad (9)$$

$$W_{\mu\nu}^{(si)}(q, p, k') = \frac{1}{2\pi} \sum_X \langle A | J_\mu(0) | k', X \rangle \langle k', X | J_\nu(0) | A \rangle \\ \times (2\pi)^4 \delta^4(p + q - k' - p_X), \quad (10)$$

where the superscript  $(si)$  denotes that it is for SIDIS. Here  $l$  and  $l'$  are the four momenta of the incoming and outgoing leptons, respectively,  $p$  is the four momentum per nucleon of the incoming target  $A$ ,  $q$  is the four momentum transfer, and  $k'$  is the four momentum of the outgoing quark. We neglect the masses and use the light-cone coordinates. The unit vectors are taken as,  $\bar{n} = (1, 0, 0, 0)$ ,  $n = (0, 1, 0, 0)$ ,  $n_{\perp 1} = (0, 0, 1, 0)$ ,  $n_{\perp 2} = (0, 0, 0, 1)$ . We choose the coordinate system such that,  $p = p^+ \bar{n}$ ,  $q = -x_B p + n Q^2 / (2x_B p^+)$ , and  $l_\perp = |\vec{l}_\perp| n_{\perp 1}$ , where  $x_B = Q^2 / 2p \cdot q$  is the Bjorken variable and  $y = p \cdot q / p \cdot l$ .

Since the azimuthal asymmetry  $\langle \cos \phi \rangle$  in the kinematic region  $k_\perp \ll Q$  is a twist-3 effect, we need to calculate  $d^2 W_{\mu\nu} / d^2 k_\perp$  up to the twist-3 level. The calculations have been carried out in Ref. [38] and we review it here for our later study of nuclear dependence.

In the unpolarized SIDIS  $e^- + A \rightarrow e^- + q + X$ , the twist-2 contribution is independent of the direction of  $\vec{k}_\perp$  and is given by,

$$\left[ \frac{d^2 W_{\mu\nu}}{d^2 k_\perp} \right]_{\text{Twist-2}} = H_{\mu\nu}^{(0)}(x_B) f_q^A(x_B, k_\perp), \quad (11)$$

where the hard part is

$$H_{\mu\nu}^{(0)}(x_B) = \frac{1}{4p \cdot q} \text{Tr} [\not{p} \gamma_\mu (x_B \not{p} + \not{q}) \gamma_\nu] = -d_{\mu\nu}, \quad (12)$$

with  $d^{\mu\nu} = g^{\mu\nu} - \bar{n}^\mu n^\nu - \bar{n}^\nu n^\mu$  as a projection tensor, and  $f_q^A(x, k_\perp)$  is TMD quark distribution function in the nucleus or nucleon,

$$f_q^A(x, k_\perp) = \int \frac{dy^- d^2 \vec{y}_\perp}{(2\pi)^3} e^{ixp^+ y^- - i\vec{k}_\perp \cdot \vec{y}_\perp} \langle A | \bar{\psi}(0) \frac{\gamma^+}{2} \mathcal{L}(0, y) \psi(y) | A \rangle, \quad (13)$$

For brevity, we now work in the covariant gauge and use  $y$  to denote four vector of coordinates  $(0, y^-, \vec{y}_\perp)$ . The gauge link  $\mathcal{L}(0, y)$  is given by,

$$\mathcal{L}(0, y) = \mathcal{L}^\dagger(-\infty, \vec{0}_\perp; 0, \vec{0}_\perp) \mathcal{L}(-\infty, \vec{y}_\perp; y^-, \vec{y}_\perp); \quad (14)$$

$$\mathcal{L}(-\infty, \vec{y}_\perp; y^-, \vec{y}_\perp) = P \exp \left( ig \int_{-\infty}^{y^-} d\xi^- A^+(\xi^-, \vec{y}_\perp) \right). \quad (15)$$

Twist-3 contributions to the semi-inclusive hadronic tensor include 3 terms which can result in  $\cos \phi$  azimuthal dependence,

$$\left[ \frac{d^2 W_{\mu\nu}}{d^2 k_\perp} \right]_{\text{Twist-3}} = \left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(0)}}{d^2 k_\perp} \right]_{\text{Twist-3}} + \left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(1,L)}}{d^2 k_\perp} \right]_{\text{Twist-3}} + \left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(1,R)}}{d^2 k_\perp} \right]_{\text{Twist-3}}, \quad (16)$$

where the number in the superscript of  $\tilde{W}_{\mu\nu}$  denotes the number of gluon(s) involved in the multiple gluon scattering with respect to which the collinear expansion is carried out; the tilde above  $W$  denotes the results after collinear expansion; and the superscript  $L$  or  $R$  denotes the left or right cut respectively with respect to the initial gluon line. The contribution  $\tilde{W}^{(0)}$  is given by,

$$\left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(0)}}{d^2 k_\perp} \right]_{\text{Twist-3}} = H_{\mu\nu, \rho}^{(0)}(x_B) \Phi_\sigma^{(0)A}(x_B, \vec{k}_\perp) d^{\rho\sigma}, \quad (17)$$

where the hard part is given by,

$$H_{\mu\nu,\rho}^{(0)}(x) = \frac{1}{4p \cdot q} \text{Tr} [\gamma_\rho \gamma_\mu (x \not{p} + \not{q}) \gamma_\nu], \quad (18)$$

and the TMD parton correlation is defined as,

$$\Phi_\sigma^{(0)A}(x, \vec{k}_\perp) = \int \frac{p^+ dy^- d^2 y_\perp}{(2\pi)^3} e^{ixp^+ y^- - i\vec{k}_\perp \cdot \vec{y}_\perp} \langle A | \bar{\psi}(0) \frac{\gamma_\sigma}{2} \mathcal{L}(0, y) \psi(y) | A \rangle. \quad (19)$$

After carrying out the trace, we obtain,

$$\left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(0)}}{d^2 k_\perp} \right]_{\text{Twist-3}} = (k_{\perp\mu} n_\nu + k_{\perp\nu} n_\mu) f_{q\perp}^A(x_B, k_\perp), \quad (20)$$

where the twist-3 parton distribution  $f_{q\perp}^A(x_B, k_\perp)$  is defined as

$$k_\perp^\alpha f_{q\perp}^A(x_B, k_\perp) = d^{\alpha\beta} \Phi_\beta^{(0)A}(x, \vec{k}_\perp). \quad (21)$$

The contribution  $\tilde{W}^{(1,R)}$  and  $\tilde{W}^{(1,L)}$  are the same in unpolarized SIDIS and are given by,

$$\frac{d^2 \tilde{W}_{\mu\nu}^{(1,R)}}{d^2 k_\perp} = \frac{d^2 \tilde{W}_{\mu\nu}^{(1,L)}}{d^2 k_\perp} = \frac{1}{4p \cdot q} \text{Tr} [\hat{h}_{\mu\nu}^{(1)\rho} \omega_\rho^{\rho'} \hat{\varphi}_{\rho'}^{(1)A}(x_B, \vec{k}_\perp)], \quad (22)$$

where  $\omega_\rho^{\rho'} \equiv g_\rho^{\rho'} - \bar{n}_\rho n^{\rho'}$ ,  $\hat{h}_{\mu\nu}^{(1)\rho} = \gamma_\mu \not{p} \gamma^\rho \not{p} \gamma_\nu$  and the matrix element

$$\hat{\varphi}_\rho^{(1)A}(x, \vec{k}_\perp) = \int \frac{p^+ dy^- d^2 y_\perp}{(2\pi)^3} e^{ixp^+ y^- - i\vec{y}_\perp \cdot \vec{k}_\perp} \langle A | \bar{\psi}(0) \mathcal{L}(0, y) D_\rho(y) \psi(y) | A \rangle, \quad (23)$$

has two independent terms contributing to the hadronic tensor

$$\left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(1)}}{d^2 k_\perp} \right]_{\text{Twist-3}} = \frac{1}{4p \cdot q} \left\{ \text{Tr} [\not{p} \hat{h}_{\mu\nu}^{(1)\rho}] k_{\perp\rho} \varphi_\perp^{(1)A}(x_B, k_\perp) + i \text{Tr} [\gamma_5 \not{p} \hat{h}_{\mu\nu}^{(1)\rho}] \epsilon_{\perp\rho\gamma} k_\perp^\gamma \tilde{\varphi}_\perp^{(1)A}(x_B, k_\perp) \right\}, \quad (24)$$

where  $\epsilon_{\perp\rho\gamma} \equiv \epsilon_{\alpha\beta\rho\gamma} \bar{n}^\alpha n^\beta$ . After carrying out the traces and make the Lorentz contraction with  $k_{\perp\rho}$  or  $\epsilon_{\perp\rho\gamma} k_\perp^\gamma$  respectively, we obtain,

$$\left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(1)}}{d^2 k_\perp} \right]_{\text{Twist-3}} = -\frac{1}{p \cdot q} (p_\mu k_{\perp\nu} + p_\nu k_{\perp\mu}) [\varphi_\perp^{(1)A}(x_B, k_\perp) - \tilde{\varphi}_\perp^{(1)A}(x_B, k_\perp)], \quad (25)$$

where the parton correlation functions are given by,

$$k_\perp^2 \varphi_\perp^{(1)A}(x, k_\perp) = n^\alpha k_\perp^\rho \varphi_{\rho\alpha}^{(1)A}(x, \vec{k}_\perp), \quad (26)$$

$$ik_\perp^2 \tilde{\varphi}_\perp^{(1)A}(x, k_\perp) = n^\alpha \epsilon_\perp^{\rho\beta} k_{\perp\beta} \tilde{\varphi}_{\rho\alpha}^{(1)A}(x, \vec{k}_\perp), \quad (27)$$

$$\varphi_{\rho\alpha}^{(1)A}(x, \vec{k}_\perp) = \int \frac{dy^- d^2 y_\perp}{(2\pi)^3} e^{ixp^+ y^- - i\vec{y}_\perp \cdot \vec{k}_\perp} \langle A | \bar{\psi}(0) \frac{\gamma_\alpha}{2} \mathcal{L}(0, y) D_\rho(y) \psi(y) | A \rangle \quad (28)$$

$$\tilde{\varphi}_{\rho\alpha}^{(1)A}(x, \vec{k}_\perp) = \int \frac{dy^- d^2 y_\perp}{(2\pi)^3} e^{ixp^+ y^- - i\vec{y}_\perp \cdot \vec{k}_\perp} \langle A | \bar{\psi}(0) \frac{\gamma_5 \gamma_\alpha}{2} \mathcal{L}(0, y) D_\rho(y) \psi(y) | A \rangle \quad (29)$$

Equation of motion relates

$$x f_{q\perp}^A(x, k_\perp) = -\varphi_\perp^{(1)A}(x, k_\perp) + \tilde{\varphi}_\perp^{(1)A}(x, k_\perp), \quad (30)$$

so we have,

$$\left[ \frac{d^2 \tilde{W}_{\mu\nu}}{d^2 k_\perp} \right]_{\text{Twist-3}} = \left[ \frac{d^2 \tilde{W}_{\mu\nu}^{(0)}}{d^2 k_\perp} + \frac{d^2 \tilde{W}_{\mu\nu}^{(1)}}{d^2 k_\perp} \right]_{\text{Twist-3}} = \frac{1}{p \cdot q} [(q_\mu + 2x_B p_\mu) k_{\perp\nu} + (q_\nu + 2x_B p_\nu) k_{\perp\mu}] f_{q\perp}^A(x_B, k_\perp), \quad (31)$$

Summing both twist-2 and twist-3 contributions and contracting with lepton tensor  $L_{\mu\nu}$ , we obtain the differential cross section as,

$$\frac{d\sigma}{dx_B dy d^2 k_\perp} = \frac{2\pi\alpha_{em}^2 e_q^2}{Q^2 y} \left\{ [1 + (1-y)^2] f_q^A(x_B, \vec{k}_\perp) - 4(2-y)\sqrt{1-y} \frac{|\vec{k}_\perp|}{Q} x_B f_{q\perp}^{(1)A}(x_B, \vec{k}_\perp) \cos \phi \right\}, \quad (32)$$

where the azimuthal angle  $\phi$  is defined by  $\cos \phi = \hat{\vec{k}}_\perp \cdot \hat{\vec{l}}_\perp$ . The azimuthal asymmetry at fixed  $k_\perp$  is given by,

$$\langle \cos \phi \rangle_{eA} = -\frac{2(2-y)\sqrt{1-y}}{1 + (1-y)^2} \frac{|\vec{k}_\perp|}{Q} \frac{x_B f_{q\perp}^A(x_B, k_\perp)}{f_q^A(x_B, k_\perp)}. \quad (33)$$

We can also calculate the transverse-momentum integrated asymmetry and obtain,

$$\langle \langle \cos \phi \rangle \rangle_{eA} = -\frac{2(2-y)\sqrt{1-y}}{1 + (1-y)^2} \frac{\int |\vec{k}_\perp| d^2 k_\perp x_B f_{q\perp}^A(x_B, k_\perp)}{f_q^A(x_B)}, \quad (34)$$

where  $f_q^A(x) = \int d^2 k_\perp f_q^A(x, k_\perp)$  is the usual quark distribution in a nucleon or nucleus.

If we consider only “free partons” with intrinsic transverse momentum, *i.e.*, setting  $g = 0$  then,  $\mathcal{L} = 1$ ,  $D_\sigma = \partial_\sigma$ , and  $x f_{q\perp}^A(x, k_\perp) = f_q^A(x, k_\perp)$ . In this case,  $\langle \cos \phi \rangle_{eA} = -2(2-y)\sqrt{1-y}/[1 + (1-y)^2] \cdot |\vec{k}_\perp|/Q$ , which is just the result obtained in Ref. [33].

We note that the above calculations apply to SIDIS of both nuclear and nucleon targets. All results are in the same form in terms of the TMD parton distributions inside a nucleus or a nucleon. The nuclear dependence of the azimuthal angle asymmetry will come from the nuclear dependence of the TMD parton distributions. In the SIDIS off a nuclear target, multiple gluon scattering between the struck quark and the target as contained in the gauge links can happen to different nucleons inside the nucleus. This will give rise to the nuclear dependence which we will discuss in the remainder of this paper.

### III. A-DEPENDENCE OF THE AZIMUTHAL ASYMMETRY

Following the same approach in the discussion of the nuclear dependence of the twist-2 TMD quark distribution  $f_q^A(x, k_\perp)$  in Ref. [24], we can also express the general parton distribution,

$$\Phi_\alpha^A(x, k_\perp) \equiv \int \frac{p^+ dy^-}{2\pi} \frac{d^2 y_\perp}{(2\pi)^2} e^{ixp^+ y^- - i\vec{k}_\perp \cdot \vec{y}_\perp} \langle A | \bar{\psi}(0) \frac{\Gamma_\alpha}{2} \mathcal{L}(0, y) \psi(y) | A \rangle, \quad (35)$$

$$= \int \frac{p^+ dy^-}{2\pi} e^{ixp^+ y^-} \langle A | \bar{\psi}(0) \frac{\Gamma_\alpha}{2} e^{\vec{W}_\perp(y^-) \cdot \vec{\nabla}_{k_\perp}} \psi(y^-) | A \rangle \delta^{(2)}(\vec{k}_\perp), \quad (36)$$

in terms of the collinear parton matrix elements involving the transport operator  $\vec{W}_\perp(y^-)$  [Eq. (2)], where  $\Gamma_\alpha$  is any gamma matrix. Expanding the exponential term of the above matrix element in powers of the transport operator  $\vec{W}_\perp(y^-)$  and assuming “maximum two-gluon correlation approximation” as in Ref. [24], one can express the nuclear TMD parton distributions in terms of a Gaussian convolution of the same TMD distributions in a nucleon,

$$\Phi_\alpha^A(x, k_\perp) \approx A \exp \left[ \frac{\Delta_{2F}}{4} \nabla_{k_\perp}^2 \right] \Phi_\alpha^N(x, k_\perp), \quad (37)$$

$$= \frac{A}{\pi \Delta_{2F}} \int d^2 \ell_\perp e^{-(\vec{k}_\perp - \vec{\ell}_\perp)^2 / \Delta_{2F}} \Phi_\alpha^N(x, \ell_\perp). \quad (38)$$

Note that in the derivation of the above result, one has considered the fact that matrix elements with odd powers of  $\vec{W}_\perp(y^-)$  vanish for  $\Gamma_\alpha = \gamma_+$  while the matrix elements with even powers of  $\vec{W}_\perp(y^-)$  vanish for  $\Gamma_\alpha = \vec{\gamma}_\perp$ .

The nuclear broadening of the twist-2 TMD quark distribution is a special case of the above general result with  $\Gamma_\alpha = \gamma_+$ . With the definition of the twist-3 parton distribution function  $f_{q\perp}^A(x, k_\perp)$  in Eqs. (19) and (21), one can

multiply both sides of Eq. (38) by  $d^{\alpha\beta}$  and obtain,

$$f_{q\perp}^A(x, k_\perp) \approx \frac{A}{\pi\Delta_{2F}} \int d^2\ell_\perp \frac{(\vec{k}_\perp \cdot \vec{\ell}_\perp)}{\vec{k}_\perp^2} e^{-(\vec{k}_\perp - \vec{\ell}_\perp)^2/\Delta_{2F}} f_{q\perp}^N(x, \ell_\perp) \quad (39)$$

$$= \frac{A}{\pi\Delta_{2F}} \left( 1 + \frac{\Delta_{2F}}{2\vec{k}_\perp^2} \vec{k}_\perp \cdot \vec{\partial}_{k_\perp} \right) \int d^2\ell_\perp e^{-(\vec{k}_\perp - \vec{\ell}_\perp)^2/\Delta_{2F}} f_{q\perp}^N(x, \ell_\perp). \quad (40)$$

Given the twist-2 TMD quark distribution function  $f_q^N(x, k_\perp)$  and the twist-3 quark distribution  $f_{q\perp}^N(x, \ell_\perp)$  in a nucleon and using the above convolution, Eqs.(3) and (40), one can then calculate the nuclear dependence of the azimuthal asymmetry  $\langle \cos \phi \rangle$  and  $\langle \langle \cos \phi \rangle \rangle$ . In general, if the convoluted twist-3 TMD quark distribution is a decreasing function of the transverse momentum in the region of interest, the second term involving a derivative in the above equation will be negative. Therefore, one would expect the azimuthal asymmetry to decrease because of the multiple scattering in nuclei.

To illustrate the nuclear dependence of the azimuthal asymmetry qualitatively, we consider an ansatz of the Gaussian distributions in  $k_\perp$  for both the twist-2 and 3 TMD quark distributions,

$$f_q^N(x, k_\perp) = \frac{1}{\pi\alpha} f_q^N(x) e^{-k_\perp^2/\alpha}, \quad (41)$$

$$f_{q\perp}^N(x, k_\perp) = \frac{1}{\pi\beta} f_{q\perp}^N(x) e^{-k_\perp^2/\beta}. \quad (42)$$

The corresponding TMD distributions in nuclei are,

$$f_q^A(x, k_\perp) \approx \frac{A}{\pi(\alpha + \Delta_{2F})} f_q^N(x) e^{-k_\perp^2/(\alpha + \Delta_{2F})}, \quad (43)$$

$$f_{q\perp}^A(x, k_\perp) \approx \frac{A\beta}{\pi(\beta + \Delta_{2F})^2} f_{q\perp}^N(x) e^{-k_\perp^2/(\beta + \Delta_{2F})}. \quad (44)$$

One can then calculate the the azimuthal asymmetry for SIDIS off both nucleon

$$\langle \cos \phi \rangle_{eN} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\alpha}{\beta} \frac{|\vec{k}_\perp|}{Q} \frac{x_B f_{q\perp}^N(x_B)}{f_q^N(x_B)} \exp\left\{-\frac{\alpha-\beta}{\alpha\beta} k_\perp^2\right\}, \quad (45)$$

and nuclear targets

$$\langle \cos \phi \rangle_{eA} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\beta(\alpha + \Delta_{2F})}{(\beta + \Delta_{2F})^2} \frac{|\vec{k}_\perp|}{Q} \frac{x_B f_{q\perp}^N(x)}{f_q^N(x)} \exp\left\{-\frac{\alpha-\beta}{(\alpha + \Delta_{2F})(\beta + \Delta_{2F})} k_\perp^2\right\}. \quad (46)$$

The nuclear modification factor for the azimuthal asymmetry is then,

$$\frac{\langle \cos \phi \rangle_{eA}}{\langle \cos \phi \rangle_{eN}} = \frac{\beta^2(\alpha + \Delta_{2F})}{\alpha(\beta + \Delta_{2F})^2} \exp\left\{\frac{(\alpha - \beta)\Delta_{2F}(\alpha + \beta + \Delta_{2F})}{\alpha\beta(\alpha + \Delta_{2F})(\beta + \Delta_{2F})} \vec{k}_\perp^2\right\}. \quad (47)$$

In the special case when  $\alpha = \beta$ , we have

$$\langle \cos \phi \rangle_{eN} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{|\vec{k}_\perp|}{Q} \frac{x_B f_{q\perp}^N(x_B)}{f_q^N(x_B)}, \quad (48)$$

$$\langle \cos \phi \rangle_{eA} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\beta}{\beta + \Delta_{2F}} \frac{|\vec{k}_\perp|}{Q} \frac{x_B f_{q\perp}^N(x)}{f_q^N(x)}, \quad (49)$$

$$\frac{\langle \cos \phi \rangle_{eA}}{\langle \cos \phi \rangle_{eN}} = \frac{\beta}{\beta + \Delta_{2F}}. \quad (50)$$

Therefore, the azimuthal asymmetry  $\langle \cos \phi \rangle_{eA}$  in deep inelastic  $eA$  scattering is suppressed compared to that in  $eN$  scattering and the suppression is inversely proportional to the total transverse momentum broadening  $\Delta_{2F}$ . The suppression is independent of the transverse momentum  $k_\perp$ . However, in general, the twist-2 and twist-3 TMD quark distributions are not necessarily the same and their Gaussian ansatz might have different widths  $\beta \neq \alpha$ . The nuclear modification factor for the azimuthal asymmetry will then have non-trivial  $k_\perp$  dependence. Shown in Fig. 1 are the nuclear modification factors for the azimuthal asymmetry when  $\beta/\alpha = 2$  and  $0.5$ , respectively, as functions of  $\Delta_{2F}/\alpha$ , at different transverse momentum  $k_\perp$ . In the case  $\beta > \alpha$ , we see that the azimuthal asymmetry is suppressed and the suppression increases with the transverse momentum  $k_\perp$ . However, when  $\beta < \alpha$ , the suppression actually decreases with increasing  $k_\perp$  and the azimuthal asymmetry could be enhanced for large enough transverse momentum  $k_\perp$ . Therefore, the nuclear modification of the azimuthal asymmetry and its transverse momentum dependence is a very sensitive probe of the twist-2 and twist-3 TMD quark distribution functions.

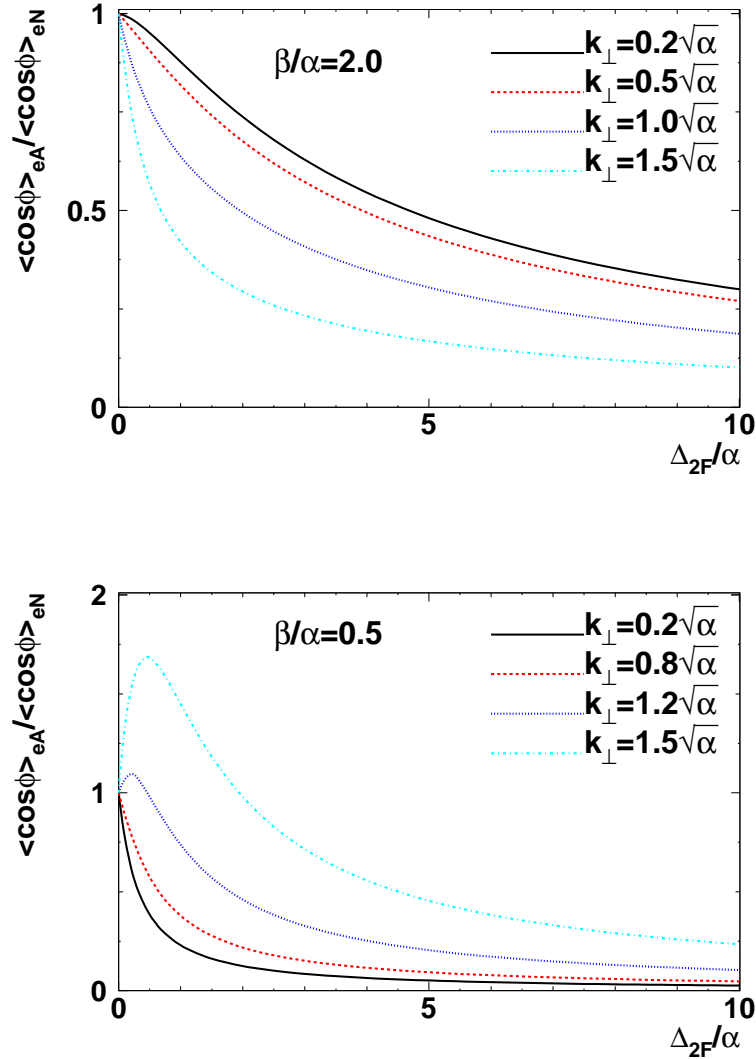


FIG. 1: (color online) Ratio  $\langle \cos \phi \rangle_{eA} / \langle \cos \phi \rangle_{eN}$  as a function of  $\Delta_{2F}$  for different  $k_\perp$  and the relative width  $\beta/\alpha$  of twist-3 and 2 TMD quark distributions.

If we integrate over the magnitude of the transverse momentum  $k_\perp$ , the averaged azimuthal asymmetry will only depend on the shape of the twist-3 TMD quark distributions,

$$\langle \langle \cos \phi \rangle \rangle_{eN} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\sqrt{\pi\beta}}{2Q} \frac{x_B f_{q\perp}^N(x)}{f_q^N(x)}, \quad (51)$$

$$\langle\langle\cos\phi\rangle\rangle_{eA} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\sqrt{\pi}\beta}{2Q\sqrt{\beta+\Delta_{2F}}} \frac{x_B f_{q\perp}^N(x)}{f_q^N(x)}, \quad (52)$$

$$\frac{\langle\langle\cos\phi\rangle\rangle_{eA}}{\langle\langle\cos\phi\rangle\rangle_{eN}} = \sqrt{\frac{\beta}{\beta+\Delta_{2F}}}. \quad (53)$$

We see again that  $\langle\langle\cos\phi\rangle\rangle_{eA}$  is suppressed compared to  $\langle\langle\cos\phi\rangle\rangle_{eN}$ , and the suppression factor is inversely proportional to the square-root of the total transverse momentum broadening  $\Delta_{2F}$ .

To take into account of the transverse momentum during the quark fragmentation process and its effect on the azimuthal asymmetry, we take another Gaussian smearing

$$d\sigma^{eN\rightarrow ehX} = \int d\sigma^{eN\rightarrow eqX} D_F^{q\rightarrow h}(z, \vec{k}_{F\perp}) dz d^2k_{F\perp} \delta^{(3)}(\vec{p}_h - z\vec{k}' - k_{F\perp}) d^3p_h, \quad (54)$$

for the  $k_{F\perp}$ -dependence in the fragmentation function  $D_F^{q\rightarrow h}(z, \vec{k}_{F\perp})$ , *i.e.*,

$$D_F^{q\rightarrow h}(z, \vec{k}_{F\perp}) = D_F^{q\rightarrow h}(z) \frac{1}{\pi\alpha_F} e^{-\vec{k}_{F\perp}^2/\alpha_F}. \quad (55)$$

Consider the case that  $\alpha = \beta$  and one has the azimuthal asymmetry for the hadron production cross section

$$\langle\cos\phi_h\rangle_{eN} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\beta z}{\beta z^2 + \alpha_F} \frac{|\vec{p}_{h\perp}|}{Q} \frac{x_B f_{q\perp}^N(x)}{f_q^N(x)}, \quad (56)$$

$$\langle\cos\phi_h\rangle_{eA} = -\frac{2(2-y)\sqrt{1-y}}{1+(1-y)^2} \frac{\beta z}{(\beta + \Delta_{2F})z^2 + \alpha_F} \frac{|\vec{p}_{h\perp}|}{Q} \frac{x_B f_{q\perp}^N(x)}{f_q^N(x)}. \quad (57)$$

We compare the above results with those obtained without fragmentation, Eqs. (48) and (49), and see that we have a clear smearing effect on the azimuthal asymmetry in both  $e^-N$  and  $e^-A$ -scatterings. The smearing factors are given by,

$$\left. \frac{\langle\cos\phi_h\rangle_{eN}}{\langle\cos\phi\rangle_{eN}} \right|_{|\vec{p}_{h\perp}|=z|\vec{k}_{\perp}|} = \frac{\beta z^2}{\beta z^2 + \alpha_F}, \quad (58)$$

$$\left. \frac{\langle\cos\phi_h\rangle_{eA}}{\langle\cos\phi\rangle_{eA}} \right|_{|\vec{p}_{h\perp}|=z|\vec{k}_{\perp}|} = \frac{(\beta + \Delta_{2F})z^2}{(\beta + \Delta_{2F})z^2 + \alpha_F}. \quad (59)$$

The suppression factor in  $eA$  compared to  $eN$  is given by,

$$\frac{\langle\cos\phi_h\rangle_{eA}}{\langle\cos\phi_h\rangle_{eN}} = \frac{\beta z^2 + \alpha_F}{(\beta + \Delta_{2F})z^2 + \alpha_F}. \quad (60)$$

After integrated over the magnitude of the transverse momentum, we have,

$$\frac{\langle\langle\cos\phi_h\rangle\rangle_{eA}}{\langle\langle\cos\phi_h\rangle\rangle_{eN}} = \sqrt{\frac{\beta z^2 + \alpha_F}{(\beta + \Delta_{2F})z^2 + \alpha_F}}. \quad (61)$$

#### IV. SUMMARY AND DISCUSSIONS

Within the generalized factorization, we have calculated the SIDIS cross sections in terms of the TMD quark distributions in a nucleon or nucleus up to twist-3. The azimuthal asymmetry  $\langle\cos\phi\rangle$  in the small transverse momentum region depends on both twist-2 and 3 TMD quark distributions. By considering nuclear broadening of both twist-2 and 3 TMD quark distributions due to multiple scattering between the struck quark and nucleons inside the nucleus,



we investigated the nuclear dependence of the azimuthal asymmetry. We found that the azimuthal asymmetry is suppressed by multiple parton scattering for most cases of the TMD quark distributions. The suppression is inversely proportional to the average squared transverse momentum broadening. The transverse momentum dependence of the suppression depends on the relative shape of the twist-2 and 3 TMD quark distributions. Using a Gaussian ansatz, we found that the suppression factor decreases with the transverse momentum if the width of the twist-2 TMD distribution is smaller than that of the twist-3 TMD distribution, while the suppression factor increases with the transverse momentum if the width of the twist-3 TMD distribution is smaller than that of the twist-2 TMD distribution. The suppression is independent of the transverse momentum if the twist-2 and 3 TMD distributions have the same width. Therefore, study of the nuclear dependence of the azimuthal asymmetry can shed light on the relative shape of the twist-2 and 3 TMD quark distributions.

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